$h \to Z\gamma$ in Type-II seesaw neutrino model

Chian-Shu Chen[2a], Chao-Qiang Geng[1,2b], Da Huang[1c], and Lu-Hsing Tsai[1d]

[1] *Department of Physics, National Tsing Hua University, Hsinchu, Taiwan*

[2] *Physics Division, National Center for Theoretical Sciences, Hsinchu, Taiwan*

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Abstract

We study the Higgs decay channel of $h \to Z\gamma$ in the type-II seesaw neutrino model. In most of the allowed parameter space in the model, the new contribution to $h \to Z\gamma$ is correlated with that to $h \to \gamma\gamma$. If the current 2σ excess of the $h \to \gamma\gamma$ rate measured by the LHC persists, the $h \to Z\gamma$ rate should be also larger than the corresponding standard model prediction. We demonstrate that the anti-correlation between $h \to \gamma\gamma$ and $h \to Z\gamma$ only exists in some special region.

[a] chianshu@phys.sinica.edu.tw

[b] geng@phys.nthu.edu.tw

[c] dahuang@phys.nthu.edu.tw

[d] lhtsai@phys.nthu.edu.tw

Current experimental results at the LHC for the Higgs search are consistent with the predictions of the Higgs boson (h) in the standard model (SM) [1, 2] except the $h \to \gamma\gamma$ decay rate measured by both ATLAS and CMS Collaborations, which is about $1.5 - 2$ times larger than the SM expectation [3]. Theoretically, models with additional charged particles in the loops are the common approaches to enhance the decay rate of $h \to \gamma\gamma$ [4]. It was pointed out that a combined analysis of $h \to \gamma\gamma$ and $h \to Z\gamma$ could provide more complete electroweak charge structure of these new physics and hence, test the feasibility of these models more precisely [5]. The Type-II seesaw mechanism [6] is a well-motivated way to generate small neutrino masses with additional charged scalars beyond the SM and its related studies on $h \to \gamma\gamma$ have been devoted in Refs. [7, 8]. The decay rate of $h \to Z\gamma$ in the Type-II seesaw model has been recently investigated in Ref. [9] and found the anti-correlated behaviors between $h \to \gamma\gamma$ and $h \to Z\gamma$. Due to the inconsistence of the formulae for the scalar contributions to $h \to Z\gamma$ used in the literature, in this paper we reanalyze the $Z\gamma$ decay rate in this model with those derived in Ref. [10]. Contrary to the previous results, we obtain a correlated relation between $h \to \gamma\gamma$ and $h \to Z\gamma$ in most of the parameter space in the Type-II seesaw model. The anti-correlation between $h \to \gamma\gamma$ and $h \to Z\gamma$ can only exist in some special case.

In the Type-II seesaw model [6], a scalar triplet Δ with its representation $(3, 2)$ under $\mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge groups is introduced, which can be expressed as

$$\Delta = \begin{pmatrix} \frac{1}{\sqrt{2}}\Delta^+ & \Delta^{++} \\ \Delta^0 & -\frac{1}{\sqrt{2}}\Delta^+ \end{pmatrix} , \tag{1}$$

leading to the Yukawa couplings

$$Y_{ab}\overline{(L_{La})^c}(i\sigma_2)\Delta(L_{Lb}) \ + \ \text{h.c.}\,, \tag{2}$$

with the Pauli matrix σ_2 and the symmetric matrix Y_{ab}. The scalar potential of the model can be in general expressed in the form

$$\begin{aligned} V(\Phi,\Delta) = & -m_\Phi^2(\Phi^\dagger\Phi) + \lambda(\Phi^\dagger\Phi)^2 + M_\Delta^2\mathrm{Tr}(\Delta^\dagger\Delta) + \lambda_1[\mathrm{Tr}(\Delta^\dagger\Delta)]^2 + \lambda_2\mathrm{Tr}\,(\Delta^\dagger\Delta)^2 \\ & +\lambda_3(\Phi^\dagger\Phi)\mathrm{Tr}(\Delta^\dagger\Delta) + \lambda_4\Phi^\dagger[\Delta^\dagger\,,\,\Delta]\Phi + \left(\frac{\mu}{\sqrt{2}}\Phi^T i\sigma_2\Delta^\dagger\Phi + \text{h.c.}\right) , \end{aligned} \tag{3}$$

where Φ is the SM Higgs doublet with the vacuum expectation value (VEV) $\langle\Phi\rangle = (0,\, v/\sqrt{2})^T$, and all parameters in the potential are taken to be real without loss of generality. Note that the potential in Eq. (3) becomes to the one in Ref. [9] after the use of the

transformations: $\lambda \to \lambda/2$, $\lambda_1 \to (\lambda_1+\lambda_2)/2$, $\lambda_2 \to -\lambda_2/2$, $\lambda_3 \to \lambda_4$, $\lambda_4 \to \lambda_5$, and $\mu \to \Lambda_6$. The neutral component Δ^0 of the triplet scalar in Eq. (1) acquires its VEV $v_\Delta = \sqrt{2}\langle\Delta^0\rangle$ through the relation

$$v_\Delta[2M_\Delta^2 + (\lambda_3-\lambda_4)v^2 + 2(\lambda_1+\lambda_2)v_\Delta^2] - \mu v^2 = 0 \,, \tag{4}$$

where M_Δ represents the mass scale of the triplet scalar. For the case of $M_\Delta \gg v$, such as the grand unification scale, the triplet VEV is naturally suppressed as $v_\Delta \approx \mu v^2/(2M_\Delta^2)$. In this scenario, the extra scalars will have no significant effects on the collider phenomena. In this paper, we concentrate on the mass scale where the triplet Δ is testable within the LHC search. In this case, we expect $M_\Delta \approx v$ so that $v_\Delta \sim \mu$. On the other hand, v_Δ is constrained to have an upper bound $v_\Delta \lesssim \mathcal{O}(1)$ GeV by the parameter $\rho \equiv m_W^2/(m_Z^2\cos^2\theta_W) = 1.004^{+0.0003}_{-0.0004}$ [11]. As a result, v_Δ comes from the nonzero coefficient μ of the last term in Eq. (3), corresponding to the breaking of lepton number symmetry. It will generate the Majorana neutrino mass at the tree level

$$(M_\nu)_{ab} = \sqrt{2}v_\Delta Y_{ab} \,, \tag{5}$$

where $a,b = e$, μ and τ. To understand the small neutrino masses, the upper bound $v_\Delta \approx 1$ GeV corresponds to a suppressed Yukawa coupling of $Y \lesssim 10^{-9}$, whereas the lower bound $v_\Delta \approx 10^{-9}$ GeV is set if $Y = \mathcal{O}(1)$. Back to the scalar sector, the mass spectra of the scalars can be solved from Eq. (3), given by

$$m_h^2 = \frac{1}{2}(M_{11}^2 + M_{22}^2 - \sqrt{(M_{11}^2 - M_{22}^2)^2 + 4M_{12}^4} \,, \tag{6}$$

$$m_{H^0}^2 = \frac{1}{2}(M_{11}^2 + M_{22}^2 + \sqrt{(M_{11}^2 - M_{22}^2)^2 + 4M_{12}^4} \,, \tag{7}$$

$$m_{A^0}^2 = \left[M_\Delta^2 + \frac{1}{2}(\lambda_3-\lambda_4)v^2 + (\lambda_1+\lambda_2)v_\Delta^2\right]\left(1+\frac{4v_\Delta^2}{v^2}\right) \,, \tag{8}$$

$$m_{H^\pm}^2 = \left[M_\Delta^2 + \frac{1}{2}\lambda_3 v^2 + (\lambda_1+\lambda_2)v_\Delta^2\right]\left(1+\frac{2v_\Delta^2}{v^2}\right) \,, \tag{9}$$

$$m_{H^{\pm\pm}}^2 = M_\Delta^2 + \frac{1}{2}\left(\lambda_3+\lambda_4\right)v^2 + \lambda_1 v_\Delta^2 \,, \tag{10}$$

where h is the SM-like Higgs, H^0 and A^0 are the CP even and odd neutral components, and H^+ and H^{++} are the singly and doubly charge mass eigenstates, respectively, while the neutral scalar mass matrix elements are

$$\begin{aligned} M_{11}^2 &= 2\lambda v^2 \,,\; M_{22}^2 = M_\Delta^2 + \frac{1}{2}(\lambda_3-\lambda_4)v^2 + 3(\lambda_1+\lambda_2)v_\Delta^2 \,, \\ M_{12}^2 &= -\frac{2v_\Delta}{v}\left[M_\Delta^2 + (\lambda_1+\lambda_2)v_\Delta^2\right] \,. \end{aligned} \tag{11}$$

The mixing angles of the singly charged and neutral scalars are approximately proportional to v_Δ/v, so the charged mass eigenstates H^+ and H^{++} nearly coincide with the weak eigenstates Δ^+ and Δ^{++}, respectively. For this reason, we will ignore the contributions from v_Δ from now on. It is also worth noticing that the trilinear couplings for the charged scalars with the SM-like Higgs h are given by

$$\mu_{hH^+H^-} = \lambda_3 v = \frac{2}{v}\left(m_{H^+}^2 - M_\Delta^2\right) , \tag{12}$$

$$\mu_{hH^{++}H^{--}} = (\lambda_3+\lambda_4) v = \frac{2}{v}\left(m_{H^{++}}^2 - M_\Delta^2\right) . \tag{13}$$

From the above relations, the deviations of the charged scalars with the triplet bare mass M_Δ clearly affect both signs and magnitudes of the corresponding trilinear couplings. In general, the mass splitting or the gauge quantum number of a scalar multiplet beyond the SM is also constrained by the oblique parameters. In the Type-II seesaw model, one can set the upper bound on the mass splitting of the triplet to be $|m_{H^{++}} - m_{H^+}| \lesssim 40$ GeV, which is insensitive to the triplet scale M_Δ [8]. The constraints for the parameters in the scalar potential can obtained from the stable conditions, given by

$$\begin{gathered}\lambda \geq 0 \ , \ \lambda_1+\lambda_2 \geq 0 \ , \ 2\lambda_1+\lambda_2 \geq 0 \ , \\ \lambda_3 \pm \lambda_4 + 2\sqrt{\lambda(\lambda_1+\lambda_2)} \geq 0 \ , \ \lambda_3 \pm \lambda_4 + 2\sqrt{\lambda(\lambda_1+\lambda_2/2)} \geq 0 \ .\end{gathered} \tag{14}$$

To ensure the perturbativity and the conditions in Eq. (14) from the electroweak to higher energy scale (e.g. Planck scale), a positive value of λ_3 is preferred if $\lambda_{1,2}$ are taken to be small [8, 9]. However, a negative value of λ_3 is still possible as long as the square roots in Eq. (14) are large enough [7]. In what follows we consider the implications of $h \to \gamma\gamma$ and $h \to Z\gamma$ in these two parameter regions.

The general formulae for scalar (s), fermion (f), and W-boson contributions to the decay rates of $h \to \gamma\gamma$ and $h \to Z\gamma$ are given by

$$\Gamma(h\to\gamma\gamma) = \frac{G_F \alpha^2 m_h^3}{128\sqrt{2}\pi^3}\left|\sum_f N_f^c Q_f^2 A_{1/2}^{\gamma\gamma}(\tau_f) + A_1^{\gamma\gamma}(\tau_W) + Q_s^2 \frac{v\mu_{hss^*}}{2m_s^2} A_0^{\gamma\gamma}(\tau_s)\right|^2 , \tag{15}$$

$$\begin{aligned}\Gamma(h\to Z\gamma) = \frac{G_F \alpha\alpha_Z m_h^3}{64\sqrt{2}\pi^3}\left(1-m_Z^2/m_h^2\right)^3 \bigg| & N_f^c \frac{Q_f(Q_R^Z+Q_L^Z)}{2} A_{1/2}^{Z\gamma}(\tau_f,\lambda_f) \\ & + Q_W Q_W^Z A_1^{Z\gamma}(\tau_W,\lambda_W) + Q_s Q_s^Z \frac{v\mu_{hss^*}}{2m_s^2} A_0^{Z\gamma}(\tau_s,\lambda_s)\bigg|^2 ,\end{aligned} \tag{16}$$

where $\alpha_Z \equiv g^2/(4\pi\cos^2\theta_W)$, N_f^c is the number of additional degrees of freedom for fermions besides the EW gauge group, μ_{hss^*} is the trilinear coupling derived from the scalar potential,

$\tau_i = m_h^2/4m_i^2$, $\lambda_i = m_Z^2/4m_i^2$, $Q_W = 1$, $Q_W^Z = \cos^2\theta_W$, $Q_{f,s}$ are the electric charges of fermions and scalars, $Q_{R,L}^Z = I_{R,L(s)}^3 - Q_{f(s)}\sin^2\theta_W$ with $I_{R,L(s)}^3$ being the third isospin components of chiral fermions (scalars), respectively. The loop functions $A_{(0,\,1/2,\,1)}^{\gamma\gamma}$ and $A_{(0,\,1/2,\,1)}^{Z\gamma}$ in Eqs. (15) and (16) are defined as

$$\begin{aligned}
A_0^{\gamma\gamma}(\tau) &= -[\tau - f(\tau)]\tau^{-2}\,, A_{1/2}^{\gamma\gamma}(\tau) = 2[\tau + (\tau-1)f(\tau)]\tau^{-2}\,,\\
A_1^{\gamma\gamma}(\tau) &= -[2\tau^2 + 3\tau + 3(2\tau-1)f(\tau)]\tau^{-2}\,,\\
A_0^{Z\gamma}(\tau,\lambda) &= I_1(\tau,\lambda)\,,\ A_{1/2}^{Z\gamma}(\tau,\lambda) = -2[I_1(\tau,\lambda) - I_2(\tau,\lambda)],\\
A_1^{Z\gamma}(\tau,\lambda) &= [2(1+2\tau)(1-\lambda) + (1-2\tau)]I_1(\tau,\lambda) - 8(1-\lambda)I_2(\tau,\lambda)\,,
\end{aligned} \tag{17}$$

where

$$\begin{aligned}
I_1(\tau,\lambda) &= -\frac{1}{(\tau-\lambda)} + \frac{1}{(\tau-\lambda)^2}[f(\tau)-f(\lambda)] + \frac{2\lambda}{(\tau-\lambda)^2}[g(\tau)-g(\lambda)]\,,\\
I_2(\tau,\lambda) &= \frac{1}{(\tau-\lambda)}[f(\tau)-f(\lambda)]\,,
\end{aligned} \tag{18}$$

with the functions $f(\tau)$ and $g(\tau)$ given by

$$f(\tau) = \begin{cases} (\sin^{-1}\sqrt{\tau})^2\,, & \tau \le 1 \\ -\frac{1}{4}[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi]^2\,, & \tau > 1 \end{cases}\,,\ g(\tau) = \begin{cases} \sqrt{\tau^{-1}-1}(\sin^{-1}\sqrt{\tau})\,, & \tau \le 1 \\ \frac{\sqrt{1-\tau^{-1}}}{2}[\log\frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi]\,, & \tau > 1\,. \end{cases} \tag{19}$$

In the SM, the W-boson contributions to $h \to \gamma\gamma$ and $h \to Z\gamma$ dominate over those from fermions, such as t-quark, and the corresponding amplitudes $A_1^{\gamma\gamma}$ and $A_1^{Z\gamma}$ are negative. The new contributions to $h \to \gamma\gamma$ or $h \to Z\gamma$ beyond the SM are usually characterized by the expressions

$$R_{\gamma\gamma(Z\gamma)} = \frac{\sigma(pp \to h)\mathrm{Br}(h \to \gamma\gamma(Z\gamma))}{\sigma_{\mathrm{SM}}(pp \to h)\mathrm{Br}_{\mathrm{SM}}(h \to \gamma\gamma(Z\gamma))}\,. \tag{20}$$

In our case, the SM-like Higgs production rates are almost the same as those for the SM since the mixing with the triplet is very small. For $h \to \gamma\gamma$, the interference between the new charged scalar and the SM contributions only depends on the sign of μ_{hss^*} since Q_s^2 is always positive. In our discussion, the trilinear couplings of H^+ and H^{++} to h are given in Eqs. (12) and (13). If μ_{hss^*} is negative (positive), then the interference with the SM one is constructive (destructive). The situation in $h \to Z\gamma$ is more complicated [10]. To determine whether the new charged scalar contribution to $h \to Z\gamma$ is constructive or destructive, we need to know the sign of not only μ_{hss^*}, but also the charge combination

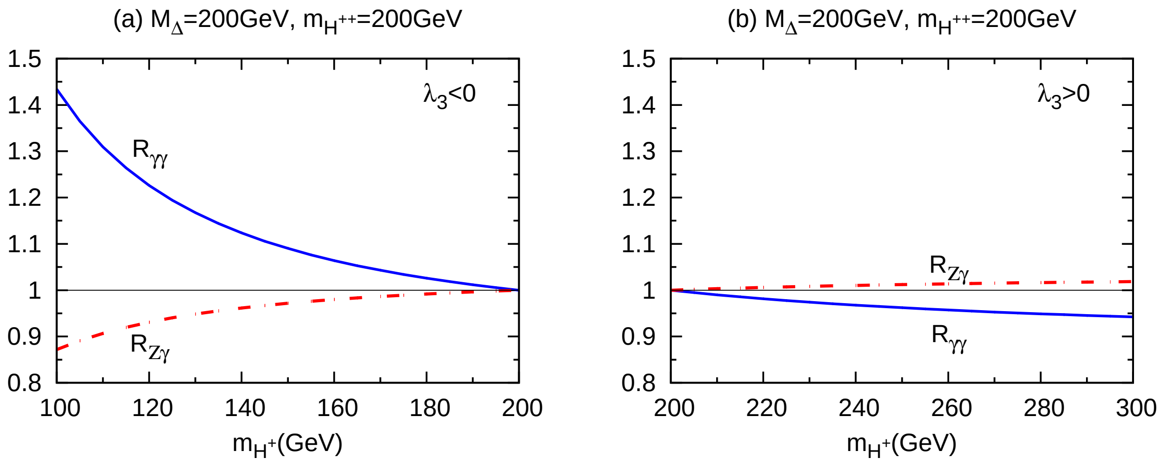

FIG. 1. $R_{\gamma\gamma}$ (solid) and $R_{Z\gamma}$ (dashed) versus m_{H^+} with $m_{H^{++}} = M_\Delta = 200$ GeV for (a) $m_{H^+} < M_\Delta$ ($\lambda_3 < 0$) and (b) $m_{H^+} > M_\Delta$ ($\lambda_3 > 0$).

$Q_s Q_s^Z = (I_3 + Y/2)(I_3 \cos^2\theta_W - Y \sin^2\theta_W/2)$. It is obvious that a larger value of I_3 yields a positive value of $Q_s Q_s^Z$, whereas $Q_s Q_s^Z$ becomes negative for a larger Y.

In the Type-II seesaw model, new contributions to $h \to \gamma\gamma$ and $h \to Z\gamma$ arise only from the loops involving with the charged scalars of $H^\pm$ and $H^{\pm\pm}$. Since the mixing between the doublet and triplet scalars is ignored, $Q_s Q_s^Z$ are negative and positive for H^+ and H^{++} as they approximately correspond to $I_3 = 0$ and 1, respectively. Clearly, for having H^+ alone, the rates of $h \to \gamma\gamma$ and $h \to Z\gamma$ are anti-correlated. We may set $m_{H^{++}} = M_\Delta$ to eliminate the contributions of H^{++} and plot with $M_\Delta = 200$GeV as presented in Fig. 1. The anti-correlated region with $m_{H^+} < M_\Delta$, corresponding to $\lambda_3 < 0$, is shown in Fig. 1a. We note that this parameter space allowed by the constraints from the vacuum stability and oblique parameters is small. On the other hand, for $m_{H^+} > M_\Delta$ with $\lambda_3 > 0$, the results are depicted in Fig. 1b. In this case, the H^+ domination is not preferred as the $h \to \gamma\gamma$ rate gets reduced, which conflicts with the current data at the LHC.

In Fig. 2, we give the related decay rates for the new contributions only from H^{++}, which is equivalent to set $m_{H^+} = M_\Delta$. In this case, the rates of $h \to \gamma\gamma$ and $h \to Z\gamma$ are correlated with each other as shown in Fig. 2. Similarly, the region with $\lambda_3 + \lambda_4 > 0$ is not preferred by the LHC results. It is important to note that for $\lambda_3 + \lambda_4 < 0$, the constraint on the mass difference between m_{H^+} and $m_{H^{++}}$ from the oblique parameters also limits the value of $m_{H^{++}}$, so that the $h \to \gamma\gamma$ rate can not be arbitrarily large. Finally, in Fig. 3 we illustrate

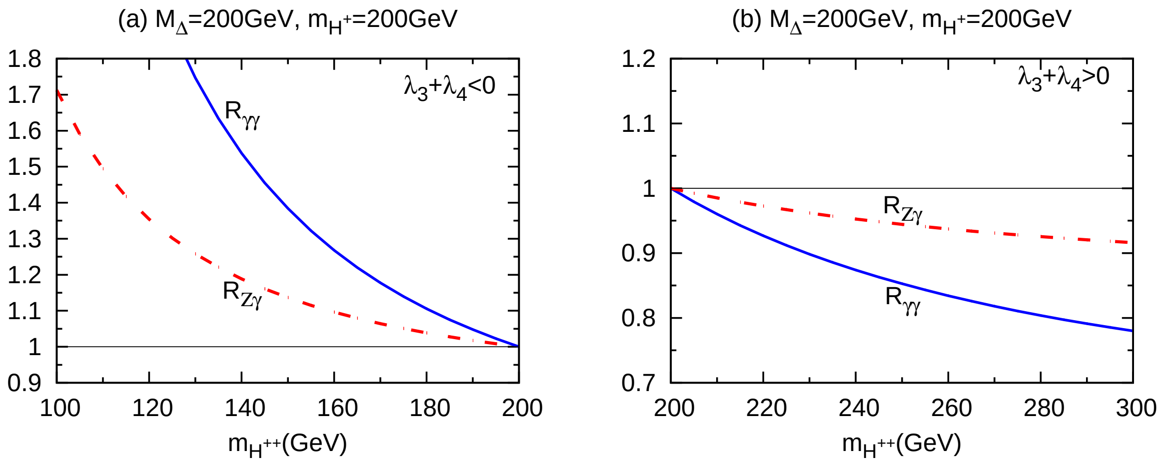

FIG. 2. $R_{\gamma\gamma}$ (solid) and $R_{Z\gamma}$ (dashed) versus $m_{H^{++}}$ with $m_{H^+} = M_\Delta = 200$ GeV for (a) $m_{H^{++}} < M_\Delta$ $(\lambda_3 + \lambda_4 < 0)$ and (b) $m_{H^{++}} > M_\Delta$ $(\lambda_3 + \lambda_4 > 0)$.

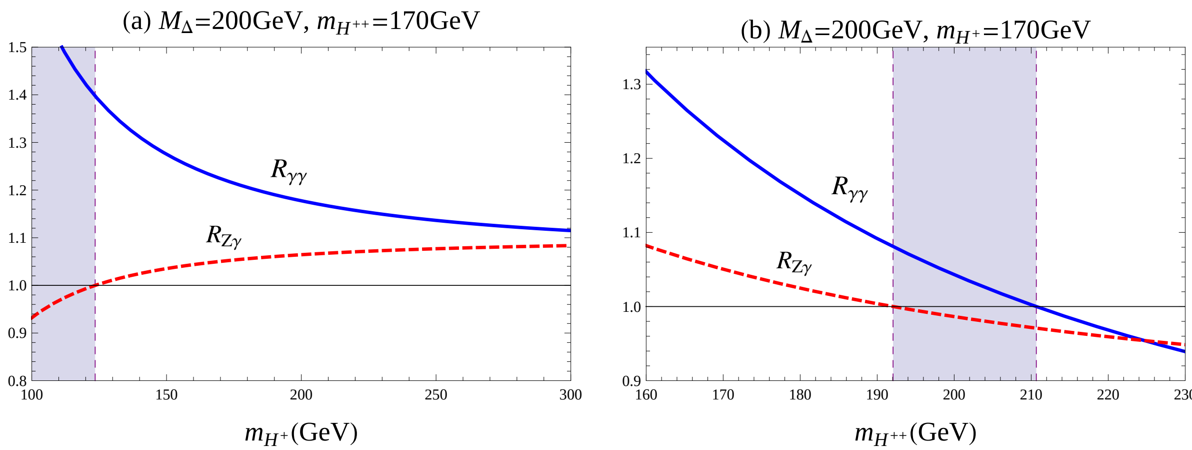

FIG. 3. $R_{\gamma\gamma}$ (solid) and $R_{Z\gamma}$ (dashed) versus (a) m_{H^+} with $m_{H^{++}} = 170$ GeV and (b) $m_{H^{++}}$ with $m_{H^+} = 170$ GeV, where $M_\Delta = 200$ GeV and the shaded areas represent the anti-correlated regions.

the general case with both H^+ and H^{++} contributions being taken into account, where we have fixed the masses of H^{++} and H^+ to be 170 GeV in Figs. 3a and 3b, respectively. It turns out that in most of the allowed parameter space, the H^{++} contributions are dominant, resulting in the positive correlation between the $h \to \gamma\gamma$ and $h \to Z\gamma$ rates, since both Q_s^2 and $Q_s Q_s^Z$ are larger than those of H^+. However, the anti-correlation can still exist if the H^+ contributions dominate over those from H^{++}. For example, one can enhance the H^+ contributions by reducing m_{H^+} and increasing $\mu_{hH^+H^+}$ simultaneously, as plotted in

Fig. 3a with $m_{H^+} \lesssim 125$ GeV. Another way is to suppress the H^{++} contributions by setting $M_\Delta \approx m_{H^{++}}$ as in the region 190 GeV $\lesssim m_{H^{++}} \lesssim$ 210 GeV in Fig. 3b.

In conclusion, we have studied the $h \to Z\gamma$ rate in the Type-II seesaw model. We have shown that the contributions to $h \to \gamma\gamma$ and $h \to Z\gamma$ from H^{++} (H^+) by itself are (anti-)correlated. On the other hand, for the general case with the existences of both H^+ and H^{++}, we have found that the deviation of the $h \to Z\gamma$ rate from the SM prediction has the same sign as the $h \to \gamma\gamma$ counterpart in most of the parameter space, whereas in some small regions with $\lambda_3 < 0$ and $m_{H^{++}} \simeq M_\Delta$, the anti-correlation between $h \to \gamma\gamma$ and $h \to Z\gamma$ appears, which could be tested in the future experiments at the LHC.

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